

# LHC 750 GeV diphoton excess and muon $(g - 2)$

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We consider implications of the diphoton excess recently observed at the LHC on the anomalous magnetic dipole moment of the muon  $(g - 2)_\mu = 2a_\mu$ , hypothesizing that the possible 750 GeV resonance is a (pseudo)scalar particle  $\tilde{\phi}$ . The  $\tilde{\phi}$ - $\gamma$ - $\gamma$  interaction implied by the diphoton events might generically contribute to  $a_\mu$  via 2-loop Barr-Zee type diagrams in a broad class of models. If  $\tilde{\phi}$  is an  $SU(2)_L$  singlet, the new contribution to  $a_\mu$  is much smaller than the current anomaly,  $\Delta a_\mu \equiv a_\mu^{\text{exp}} - a_\mu^{\text{SM}} \approx (30 \pm 10) \times 10^{-10}$ , since the scalar can complete the Barr-Zee diagrams only through its mixing with the Standard Model Higgs boson. If  $\tilde{\phi}$  belongs to an  $SU(2)_L$  doublet in an extended Higgs sector, then by contrast,  $\Delta a_\mu$  can be easily accommodated with the aid of an enhanced Yukawa coupling of  $\tilde{\phi}$  to the muon such as in the Type-II or -X two Higgs doublet model.

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Recently, both the ATLAS and the CMS collaborations at the LHC observed a possible resonance around 750 GeV in the diphoton mass distribution from the dataset of  $pp$  collisions at  $\sqrt{s} = 13$  TeV [1, 2]. The significance of deviation reported by ATLAS (CMS) is  $3.9$  ( $2.6$ )  $\sigma$  out of the  $3.2$  ( $2.6$ )  $\text{fb}^{-1}$  sample if the look-elsewhere effect is not taken into account. The measured excess in the cross section is

$$\sigma(pp \rightarrow \gamma\gamma) = \begin{cases} (10 \pm 3) \text{ fb} & \text{ATLAS} \\ (6 \pm 3) \text{ fb} & \text{CMS.} \end{cases} \quad (1)$$

Clearly, more data are called for to confirm or exclude this intriguing hint at new physics.

In the meantime, a large number of works have already appeared to put forth diverse theories on a yet-unknown resonance [3–15]. From this sharp burst of endeavours, an outstanding property has emerged which is common in the majority of the phenomenological models: the resonance candidate is required to have a rather strong interaction with a pair of photons to reproduce the experimentally preferred event rate [3]. For instance, the same type of triangle diagrams as for the Standard Model (SM) Higgs decay into two photons would not suffice. This has naturally led us to think of a possible connection between the diphoton excess and another popular observable susceptible to new electromagnetic interactions, i.e. the anomalous magnetic moment of the muon  $(g - 2)_\mu = 2a_\mu$ . It is well known to show a long-standing deviation of about  $3\sigma$  from the SM prediction [16]:

$$\Delta a_\mu \equiv a_\mu^{\text{exp}} - a_\mu^{\text{SM}} = (29.0 \pm 9.0 \quad \text{to} \quad 33.5 \pm 8.2) \times 10^{-10}. \quad (2)$$

Our aim shall then be to see whether or not this discrepancy can be ameliorated by generic properties of the newly introduced resonance.

To this end, we shall focus on a broad class of models in which the 750 GeV resonance is a spin-zero boson. The pivotal point here is that the (pseudo)scalar-diphoton vertex necessary to produce the excess could generically contribute to the Barr-Zee type diagrams [17] for  $a_\mu$  [18–23], provided that the resonance couples to the muon. For this Yukawa coupling, two mechanisms are conceivable: indirect or direct. The former is presumably most generic in the sense that it would allow the boson to interact with muons even if the muon has a Yukawa coupling only with the SM Higgs doublet. Prime examples in which this is the case would be singlet-extensions of the SM (see e.g. [8, 10–12]). In this class of models, the gauge symmetry forbids a renormalizable coupling between a new  $SU(2)_L$ -singlet scalar and muons. It is well known nevertheless that a heavy mass eigenstate can couple to muons through the mixing between the singlet and the SM Higgs boson. This leads however to the drawback that the coupling is suppressed by the mixing angle. The latter, direct mechanism is more straightforward. One may simply introduce an additional  $SU(2)_L$ -doublet which can form a Yukawa coupling with a muon pair. Obvious examples include two Higgs doublet models (2HDMs) (see e.g. [5, 8–11, 13, 14]). An important feature of this class of models is that the resonance-muon coupling can be stronger than the SM muon Yukawa coupling depending on the structure of the Higgs-Yukawa sector, of which one might take advantage to explain  $\Delta a_\mu$ . In 2HDMs for instance, this would amount to playing with the Higgs mixing angles and the Yukawa “types”.

In what follows, we evaluate the Barr-Zee type  $a_\mu$  diagrams induced by the resonance-diphoton effective vertex embedded in them. We consider elementary cases where the effective vertex is dominated by a single particle circulating in the loop, which allows us to obtain a simple relation between  $\Delta a_\mu$  and the diphoton decay amplitude. By using the resonance-diphoton coupling strength preferred by the ATLAS and CMS data, we predict

the range of  $\Delta a_\mu$ . Furthermore, we comment on a popular case where the effective vertex arises from vector-like fermions which can form multiple states in the loop.

Shortly after the announcement from the LHC, a paper appeared which included qualitative discussion of the Barr-Zee type contributions to  $a_\mu$  mediated by a 750 GeV resonance [6]. In our work, we perform a more quantitative analysis based on concrete prescriptions. Another paper included the Barr-Zee graphs although they did not play a significant role in the results [9]. Apart from the Barr-Zee diagrams, other types of corrections to  $a_\mu$  might arise which depend on the details of each model for the diphoton excess [4, 7, 8].

We begin the analysis by introducing an effective Lagrangian [24] to describe the decay of a 750 GeV (pseudo)scalar  $\tilde{\phi}$  into two photons:

$$\mathcal{L}_{\text{eff}} = \tilde{c}_\gamma \frac{\alpha}{\pi v} \tilde{\phi} F_{\mu\nu} \tilde{F}^{\mu\nu}, \quad (3)$$

where the electromagnetic dual field strength tensor is given by  $\tilde{F}^{\mu\nu} = \epsilon^{\mu\nu\rho\sigma} F_{\rho\sigma}/2$  and  $v \approx 246$  GeV is the vacuum expectation value (VEV) of the SM Higgs field. One can express the effective coupling strength  $|\tilde{c}_\gamma|$  needed to fit the LHC diphoton excess in the form

$$|\tilde{c}_\gamma^{\text{LHC}}| \simeq 5.0 \times \left( \frac{\Gamma_{\gamma\gamma}/m_{\tilde{\phi}}}{1.0 \times 10^{-4}} \right)^{1/2}, \quad (4)$$

in terms of the decay width  $\Gamma_{\gamma\gamma} \equiv \Gamma(\tilde{\phi} \rightarrow \gamma\gamma)$ . This relation allows us to estimate a viable range of  $|\tilde{c}_\gamma|$  from that of  $\Gamma_{\gamma\gamma}$ ,

$$1.1 \times 10^{-6} \lesssim \Gamma_{\gamma\gamma}/m_{\tilde{\phi}} \lesssim 2 \times 10^{-3}, \quad (5)$$

determined from its dependence on  $\Gamma_{gg} \equiv \Gamma(\tilde{\phi} \rightarrow gg)$  as well as  $\sigma(pp \rightarrow \tilde{\phi} \rightarrow \gamma\gamma)$  [3]. The lower limit arises from the condition that  $\tilde{\phi} \rightarrow \gamma\gamma$  and  $\tilde{\phi} \rightarrow gg$  saturate the total width  $\Gamma_{\text{tot}}$  while the upper limit applies when the  $\tilde{\phi}$  production is dominated by photon fusion.

Values of  $|\tilde{c}_\gamma^{\text{LHC}}|$  from (4) and (5) are much larger than the size of  $c_\gamma$  which would result if the 750 GeV scalar had only SM-like interactions. More concretely, this coupling would read

$$c_\gamma^{\text{SML}} = c_\gamma^{\text{SML}}(f) + c_\gamma^{\text{SML}}(V), \quad (6)$$

where the fermion and the vector-boson loop contributions,

$$c_\gamma^{\text{SML}}(f) = \frac{N(r_t)Q_t^2}{6} A_f(\tau_t), \quad (7a)$$

$$c_\gamma^{\text{SML}}(V) = -\frac{7}{8} A_v(\tau_W), \quad (7b)$$

would result from the SM-like top- and  $W$ -loops. Here,  $N(r_t) = 3$  is the number of top quarks with different colours,  $Q_t = 2/3$  is the top quark charge, the loop functions  $A_f(\tau)$  and  $A_v(\tau)$  are given in [24], and  $\tau_t = m_\phi^2/4m_t^2$ ,  $\tau_W = m_\phi^2/4m_W^2$ . The numerical size of (6) would then be<sup>1</sup>

$$|c_\gamma^{\text{SML}}| \simeq 0.087, \quad (8)$$

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<sup>1</sup> We consider only the scalar case. For the pseudoscalar case, the vector boson contribution should be omitted.

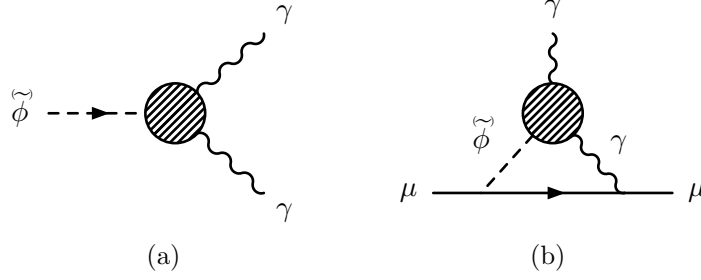


FIG. 1. Generic diagram for  $\tilde{\phi} \rightarrow \gamma\gamma$  (a) and Barr-Zee type diagram for  $a_\mu$  (b).

which is two orders of magnitude smaller than the typical size of  $\tilde{c}_\gamma^{\text{LHC}}$  from (4).

This requires contributions to  $\tilde{c}_\gamma$  much bigger than  $c_\gamma^{\text{SML}}$ , presumably arising from certain underlying physics. We sketch a generic diagram for this in Fig. 1(a). This graph can be embedded in Fig. 1(b) thereby generating a contribution to  $a_\mu$ , provided that  $\tilde{\phi}$  couples to muons.

In view of the large  $|\tilde{c}_\gamma|$  from (4), it might be expected to induce a sizeable  $\Delta a_\mu$ . An obstacle is however that the effective operator in (3) cannot be used for a direct calculation of the diagram in Fig. 1(b). This is mainly due to the difference between the kinematics involved in the  $\tilde{\phi} \rightarrow \gamma\gamma$  and  $\Delta a_\mu$  calculations. For example, the photons in Fig. 1(a) are highly energetic while the external photon in Fig. 1(b) is very soft. A naive application of (3) would lead to both ultraviolet and infrared divergences in  $\Delta a_\mu$ .

To circumvent these problems, we shall assume that the  $\tilde{\phi}$ - $\gamma$ - $\gamma$  vertex originates from loops of heavy particles. Moreover, we shall mainly focus on cases where the effective vertex is dominated by a one-loop contribution involving a single particle, to make the  $\Delta a_\mu$  predictions as model-independent as possible. As we will see, relaxing this single-particle dominance would lead to similar conclusions. We shall consider three types of particles that could appear inside the loop in Fig. 1(a): fermion ( $f$ ), vector ( $V$ ), and scalar ( $S$ ). Their interaction Lagrangian might read

$$\mathcal{L} = -\xi_\phi^f \frac{m_f}{v} \phi \bar{f} f + \rho_\phi^V \frac{2m_V^2}{v} \phi V_\mu^\dagger V^\mu - \lambda_\phi^S v \phi S^\dagger S + i\xi_\phi^f \frac{m_f}{v} \tilde{\phi} \bar{f} \gamma_5 f. \quad (9)$$

The individual contributions to  $\tilde{c}_\gamma$  can then be written in the forms [24],

$$c_\gamma(f) = \frac{N(r_f) Q_f^2 \xi_\phi^f}{6} A_f(\tau_f), \quad (10a)$$

$$c_\gamma(V) = -\frac{7N(r_V) Q_V^2 \rho_\phi^V}{8} A_v(\tau_V), \quad (10b)$$

$$c_\gamma(S) = \frac{N(r_S) Q_S^2 \lambda_\phi^S v^2}{48m_S^2} A_s(\tau_S), \quad (10c)$$

$$\tilde{c}_\gamma(f) = -\frac{N(r_f) Q_f^2 \xi_\phi^f}{4} A_a(\tau_f), \quad (10d)$$

where  $\tau_i = m_\phi^2/4m_i^2$  with  $i = f, V, S$ , and the loop functions  $A_{f,v,s}(\tau)$  and  $A_a(\tau)$  are given in [24, 25], respectively. For each of these four types of vertices, one can evaluate the

corresponding two-loop Barr-Zee graph in Fig. 1(b) to obtain the generic formulae [18, 22],

$$\Delta a_\mu(f) = \frac{\alpha m_\mu^2}{4\pi^3 v^2} N(r_f) Q_f^2 \xi_\phi^f \xi_\phi^\mu \mathcal{F}_f(z_{f\phi}), \quad (11a)$$

$$\Delta a_\mu(V) = \frac{\alpha m_\mu^2}{8\pi^3 v^2} N(r_V) Q_V^2 \rho_\phi^V \xi_\phi^\mu \mathcal{F}_v(z_{V\phi}), \quad (11b)$$

$$\Delta a_\mu(S) = \frac{\alpha m_\mu^2}{8\pi^3 m_S^2} N(r_S) Q_S^2 \lambda_\phi^S \xi_\phi^\mu \mathcal{F}_s(z_{S\phi}), \quad (11c)$$

$$\Delta \tilde{a}_\mu(f) = \frac{\alpha m_\mu^2}{4\pi^3 v^2} N(r_f) Q_f^2 \xi_\phi^f \xi_\phi^\mu \tilde{\mathcal{F}}_f(z_{f\tilde{\phi}}), \quad (11d)$$

where  $z_{i\tilde{\phi}} = m_i^2/m_{\tilde{\phi}}^2$  with  $i = f, V, S$ , and the loop functions are given by

$$\mathcal{F}_f(z) = \frac{z}{2} \int_0^1 dx \frac{2x(1-x) - 1}{z - x(1-x)} \log \frac{z}{x(1-x)}, \quad (12a)$$

$$\mathcal{F}_v(z) = \frac{1}{2} \int_0^1 dx \frac{x(12x^2 - 3x + 10)z - x(1-x)}{z - x(1-x)} \log \frac{z}{x(1-x)}, \quad (12b)$$

$$\mathcal{F}_s(z) = \frac{z}{2} \int_0^1 dx \frac{x(x-1)}{z - x(1-x)} \log \frac{z}{x(1-x)}, \quad (12c)$$

$$\tilde{\mathcal{F}}_f(z) = \frac{z}{2} \int_0^1 dx \frac{1}{z - x(1-x)} \log \frac{z}{x(1-x)}. \quad (12d)$$

With the above ingredients, we are ready to consider the first class of models, i.e. those with an additional scalar field  $\Phi$  (complex or real) which is singlet under the SM gauge group. Even though the gauge symmetry forbids a direct  $\Phi$ - $\mu$ - $\mu$  coupling, a renormalizable interaction of the form,

$$\lambda_{H\Phi} H^\dagger H \Phi^\dagger \Phi, \quad (13)$$

can mix  $\phi$ , a real degree of freedom out of  $\Phi$ , with  $h$ , the SM Higgs from the Higgs doublet  $H$ , if both  $H$  and  $\Phi$  acquire VEVs. Barring  $CP$ -violation,  $\phi$  here should be a  $CP$ -even scalar. The lighter and the heavier mass eigenstates  $H_1$  and  $H_2$  can then be identified with the 125 GeV Higgs boson and the 750 GeV resonance, respectively. We let  $R(\theta)$  denote the mixing matrix

$$R(\theta) = \begin{pmatrix} \cos \theta & -\sin \theta \\ \sin \theta & \cos \theta \end{pmatrix}, \quad (14)$$

so that

$$\begin{pmatrix} h \\ \phi \end{pmatrix} = R(-\alpha) \begin{pmatrix} H_1 \\ H_2 \end{pmatrix}. \quad (15)$$

In terms of the mixing angle  $\alpha$ , the three pieces of couplings for  $H_2 \rightarrow \gamma\gamma$  are written as

$$c_\gamma^{H_2}(f) = c_\alpha c_\gamma(f) + s_\alpha c_\gamma^{\text{SML}}(f), \quad (16a)$$

$$c_\gamma^{H_2}(V) = c_\alpha c_\gamma(V) + s_\alpha c_\gamma^{\text{SML}}(V), \quad (16b)$$

$$c_\gamma^{H_2}(S) = c_\alpha c_\gamma(S), \quad (16c)$$

where the  $\phi$ -components and the SM-like components are from (10) and (7), respectively.

Since the SM result of the Higgs decay into two photons agrees with the LHC data [26],  $R$  should be close to a unit matrix with  $|\alpha| \ll 1$ . Given that the SM-like contribution  $c_\gamma^{\text{SM}}$  has a small size  $\sim 0.087$  [see (8)], we can safely neglect the terms suppressed by  $s_\alpha$  in (16). For each particle type of fermion, vector, and scalar, dominating the  $\phi$ - $\gamma$ - $\gamma$  loop, we then get a ratio of  $\Delta a_\mu$  to  $c_\gamma^{H_2}$  in one of the simple forms,

$$\frac{\Delta a_\mu(f)}{c_\gamma^{H_2}(f)} \simeq -\frac{3\alpha s_\alpha m_\mu^2}{2\pi^3 v^2} \frac{\mathcal{F}_f(z_{fH_1}) - \mathcal{F}_f(z_{fH_2})}{A_f(\tau_f)}, \quad (17a)$$

$$\frac{\Delta a_\mu(V)}{c_\gamma^{H_2}(V)} \simeq \frac{\alpha s_\alpha m_\mu^2}{7\pi^3 v^2} \frac{\mathcal{F}_v(z_{VH_1}) - \mathcal{F}_v(z_{VH_2})}{A_v(\tau_V)}, \quad (17b)$$

$$\frac{\Delta a_\mu(S)}{c_\gamma^{H_2}(S)} \simeq -\frac{6\alpha s_\alpha m_\mu^2}{\pi^3 v^2} \frac{\mathcal{F}_s(z_{SH_1}) - \mathcal{F}_s(z_{SH_2})}{A_s(\tau_S)}. \quad (17c)$$

We can see that many parameters have been cancelled out of the numerator and denominator. In Fig. 2, we plot the above ratios as functions of  $m_X$  ( $X = f, V, S$ ), the particle mass in the  $\phi$ - $\gamma$ - $\gamma$  loop. We fixed  $s_\alpha = 0.1$ , a representative small mixing which is still allowed [27]. It is interesting to find that the ratios are similar to one another especially in the decoupling regime. Given that  $|\tilde{c}_\gamma^{\text{LHC}}|$  in (4) is around 5 we can *predict*  $\Delta a_\mu$  to be around a few  $\times 10^{-11}$  or less. This prediction is however too small to explain the deviation in (2).

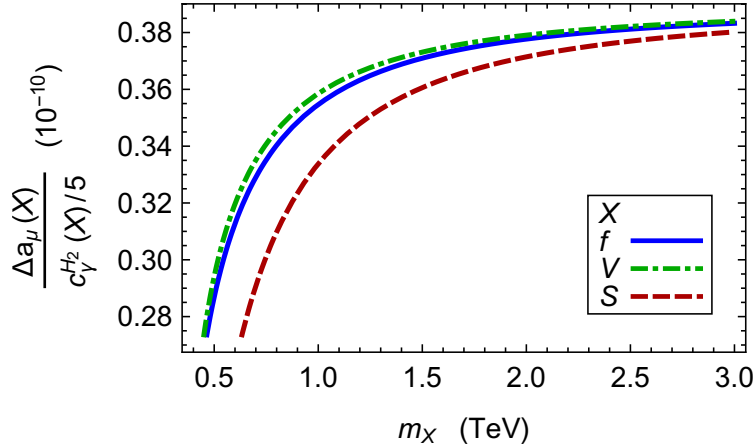


FIG. 2. The ratio  $\Delta a_\mu(X)/c_\gamma^{H_2}(X)$  as a function of  $m_X$ , mass of the particle  $X$  in the  $\phi$ - $\gamma$ - $\gamma$  loop, for  $X = f, V, S$ , standing for fermion (solid), vector (dash-dotted), or scalar (dashed), respectively.

Now we turn to the second class of models, in which the Higgs sector includes new scalars transforming non-trivially under the  $SU(2)_L$  group. Let us consider a 2HDM as the simplest example. Our results could be easily extended to more complicated models with additional doublets. In 2HDMs, a  $Z_2$  symmetry is often imposed to prevent dangerous flavour changing neutral currents mediated by Higgs at tree-level [28]. Depending on the assignment of the  $Z_2$  parity, there are four different types of naturally flavour-conserving (NFC) models [29], named, Type-I, -II, -X, and -Y (see e.g. [30]). In the Type-II and Type-X models, the Yukawa coupling of the heavier (pseudo)scalar  $H(A)$  to leptons can be enhanced by the factor  $\xi_{H(A)}$ . Both these factors become  $\tan \beta$  in the alignment limit with  $\sin(\beta - \alpha) = 1$ , which we shall adopt for the lightest state to have SM-like properties [31]. This also makes

it simpler to consider heavy Higgs decays by suppressing the  $H$  couplings to vector bosons and to  $hh$ . If  $H$  or  $A$  is identified with the 750 GeV resonance, the ratio  $\Delta\tilde{a}_\mu/\tilde{c}_\gamma$  can be written approximately as

$$\frac{\Delta a_\mu(f)}{c_\gamma(f)} \simeq \frac{3\alpha\xi_H^\mu m_\mu^2}{2\pi^3 v^2} \frac{\mathcal{F}_f(z_{fH})}{A_f(\tau_f)}, \quad (18a)$$

$$\frac{\Delta a_\mu(V)}{c_\gamma(V)} \simeq -\frac{\alpha\xi_H^\mu m_\mu^2}{7\pi^3 v^2} \frac{\mathcal{F}_v(z_{VH})}{A_v(\tau_V)}, \quad (18b)$$

$$\frac{\Delta a_\mu(S)}{c_\gamma(S)} \simeq \frac{6\alpha\xi_H^\mu m_\mu^2}{\pi^3 v^2} \frac{\mathcal{F}_s(z_{SH})}{A_s(\tau_S)}, \quad (18c)$$

$$\frac{\Delta\tilde{a}_\mu(f)}{\tilde{c}_\gamma(f)} \simeq -\frac{\alpha\xi_A^\mu m_\mu^2}{\pi^3 v^2} \frac{\tilde{\mathcal{F}}_f(z_{fA})}{A_a(\tau_f)}, \quad (18d)$$

keeping only the  $\tan\beta$ -enhanced terms. Fig. 3 shows the above ratios as functions of  $m_X$ , the particle mass in the  $\tilde{\phi}$ - $\gamma$ - $\gamma$  loop. We took  $\xi_H^\mu = \xi_A^\mu = \tan\beta = 10$ . It is remarkable that the  $\Delta a_\mu$  in (2) can be easily explained with  $|\tilde{c}_\gamma^{\text{LHC}}| \sim 5$  in (4) if  $\tilde{c}_\gamma < 0$ . For this, one could make the sign of each piece of  $\tilde{c}_\gamma$  negative by choosing appropriate signs of the coupling constants appearing in (9). We also note that the ratios have the same sign and similar magnitudes. This implies enough possibility that the ratio of total  $\Delta a_\mu$  to total  $\tilde{c}_\gamma$  remains somewhere among the curves in Fig. 3 even in the presence of several simultaneous contributions from  $f$ ,  $V$ , and  $S$ , with comparable sizes.

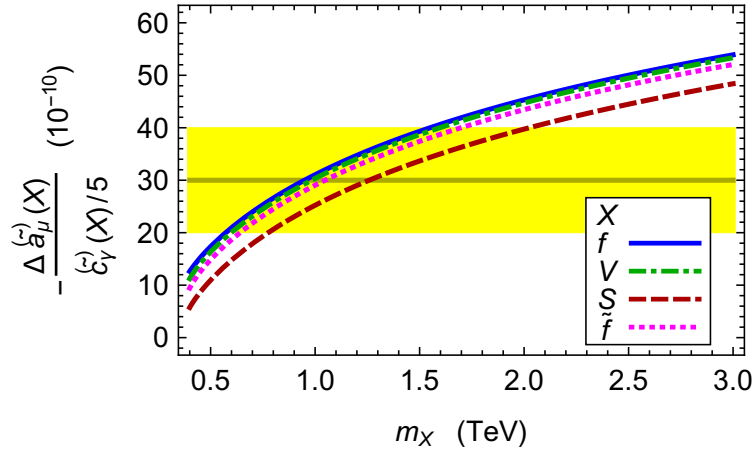


FIG. 3. The ratio  $\Delta a_\mu(X)/c_\gamma(X)$  as a function of  $m_X$ , mass of the particle  $X$  in the  $\phi$ - $\gamma$ - $\gamma$  loop, for  $X = f, V, S$ , standing for fermion (solid), vector (dash-dotted), or scalar (dashed), respectively, as well as  $\Delta\tilde{a}_\mu(f)/\tilde{c}_\gamma(f)$  as a function of  $m_f$  (dotted). The horizontal band depicts the experimentally preferred range of  $\Delta a_\mu$ .

Since we rely on the enhancement of the  $\tilde{\phi}$ - $\mu$ - $\mu$  coupling, remarks on the decays of  $\tilde{\phi}$  into lepton pairs are in order [5, 23]. The expression for  $\Gamma_{\ell\ell} \equiv \Gamma(\tilde{\phi} \rightarrow \ell^+\ell^-)$  for a lepton  $\ell$ ,

$$\frac{\Gamma_{\ell\ell}}{\Gamma_{\gamma\gamma}} \simeq \frac{\pi^2 m_\ell^2 \xi_\phi^{\ell 2}}{2\alpha^2 m_\phi^2 \tilde{c}_\gamma^2}, \quad (19)$$

can be put in the semi-numerical form,

$$\left[ \frac{\Gamma_{\mu\mu}/\Gamma_{\gamma\gamma}}{\Gamma_{\tau\tau}/\Gamma_{\gamma\gamma}} \right] \simeq \left[ \frac{7.4 \times 10^{-3}}{2.1} \right] \left( \frac{5}{\tilde{c}_\gamma} \right)^2 \left( \frac{\xi_\phi^{\mu,\tau}}{10} \right)^2. \quad (20)$$

From this, it is obvious that the bound  $\Gamma_{\mu\mu}/\Gamma_{\gamma\gamma} \lesssim 0.6$  from [3] is fulfilled by the parameters chosen above. Even if one further assumes that the  $\tilde{\phi}$ - $\tau$ - $\tau$  coupling is enhanced by the same factor as in NFC 2HDMs, the bound  $\Gamma_{\tau\tau}/\Gamma_{\gamma\gamma} \lesssim 6$  from [3] is still satisfied albeit with a less margin. One can then multiply (20) by (1) to estimate the dilepton production cross sections through the resonance at 13 TeV,

$$\left[ \frac{\sigma(pp \rightarrow \mu^+ \mu^-)}{\sigma(pp \rightarrow \tau^+ \tau^-)} \right] \simeq \left[ \frac{0.059 \text{ fb}}{17 \text{ fb}} \right] \left( \frac{5}{\tilde{c}_\gamma} \right)^2 \left( \frac{\xi_\phi^{\mu,\tau}}{10} \right)^2, \quad (21)$$

where we have taken an average of the two diphoton cross sections. These channels might be handles to cross check the present proposal at future runs of the LHC. Notice that the 13 TeV tau pair cross section might already exceed 12 fb, the 95% confidence level upper limit from the 8 TeV run [32]. If heavy Higgs searches in this channel rule out the tau pair production rate suggested by  $\Delta a_\mu$  in the near future, then one might resort to a non-NFC scenario where  $\xi_\phi^\mu \neq \xi_\phi^\tau$ .

Depending on the model, the decay  $\tilde{\phi} \rightarrow b\bar{b}$  might also be enhanced as is the case e.g. in the Type-II 2HDM, in contrast to the Type-X model where the  $b\bar{b}$  mode would be suppressed. To compensate for the suppression of  $\text{BR}(\tilde{\phi} \rightarrow \gamma\gamma)$ , the former model would therefore require a higher production rate of  $pp \rightarrow \tilde{\phi}$  than the latter. We shall leave the construction of a resonance production mechanism out of the scope of this work, as our discussion is independent of a concrete realization thereof. In such a model, it would also be of interest to look for heavy Higgs decays into  $b\bar{b}$ , whose event rate would be roughly 10 times that of the  $\tau^+\tau^-$  channel shown in (21).

In the Type-II model,  $b \rightarrow s\gamma$  places a lower limit on the charged Higgs mass  $m_{H^\pm}$  around 480 GeV at the 95% confidence level [33]. This is easy to satisfy without spoiling the  $\rho$  parameter by assuming that  $m_{H^\pm} \sim m_A \sim 750$  GeV. To probe an effect from such a heavy charged Higgs on  $b \rightarrow s\gamma$ , one would need a better experimental precision than would be available at a super flavour factory as well as more accurate theory predictions [see e.g. Fig. 7(b) of [34]].

The rather large modulus of  $\tilde{c}_\gamma$  for a high enough rate of  $\tilde{\phi} \rightarrow \gamma\gamma$  is a challenge to all weakly-coupled models (see e.g. [3]). This is even more the case if the same type of coupling is to make a sufficient contribution to  $a_\mu$  [6]. For instance,  $c_\gamma(f) \sim -5$  was shown above to fit the central value of  $\Delta a_\mu$ . This would require  $|N(r_f)Q_f^2\xi_\phi^f| \sim 29$  when  $m_f \sim 1$  TeV for instance, which may still be considered to lie within the boundary of perturbativity.

A popular way to model massive charged fermions which couple to heavier Higgses is to make them vector-like [5, 8, 9, 11, 13]. For instance, a single vector-like “generation” might consist of the following fermions:

$$l \text{ (}\mathbf{2}, Q_f + 1/2\text{)}; l^c \text{ (}\mathbf{2}, -Q_f - 1/2\text{)}; e \text{ (}\mathbf{1}, Q_f\text{)}; e^c \text{ (}\mathbf{1}, -Q_f\text{)}; \quad (22)$$

where each left-handed Weyl spinor is followed by its  $SU(2)_L$  representation and hypercharge enclosed in parentheses. Suppose that these new fermions have  $Z_2$  parities such that they



can couple to one of the Higgs doublets  $\Phi_i$  with  $i = 1$  or  $2$ . This would allow the following terms in the Lagrangian:

$$-\mathcal{L} = y\Phi_i^\dagger l e^c + y^c \Phi_i^T l^c e + m_l l l^c + m_e e e^c + \text{h.c.} \quad (23)$$

With the VEV of  $\Phi_i$  taken into account, the above Weyl fermions comprise two Dirac mass eigenstates, each of which can be regarded as  $f$  appearing in (9). One can obtain the coefficients  $\xi_\phi^f$  therein in terms of  $y$ ,  $y^c$ , the mass eigenvalues and mixing matrices of the vector-like fermions, as well as the Higgs mixing angles. For simplicity, we shall assume all the parameters to be real-valued in (23). Using mass insertion approximation, one can then express the sum of contributions to  $\tilde{c}_\gamma$  from the pair of  $f$  in the forms,

$$c_\gamma(f) \simeq -\frac{N(r_f)Q_f^2}{3} \frac{R_{i1}^\beta R_{i1}^\alpha v^2}{m_\phi^2} \left[ (y + y^c)^2 \tau^2 A'_f(\tau) + 2yy^c \tau A_f(\tau) \right], \quad (24a)$$

$$\tilde{c}_\gamma(f) \simeq \frac{N(r_f)Q_f^2}{2} \frac{R_{i1}^\beta R_{i2}^\beta v^2}{m_\phi^2} (y^{c2} - y^2) \tau^2 A'_a(\tau), \quad (24d)$$

where  $R^\alpha$  and  $R^\beta$  are respectively the abbreviations of the Higgs mixing matrices  $R(\alpha)$  and  $R(\beta)$  as shown in (14), and  $\tau = m_\phi^2/2(m_l^2 + m_e^2)$ . A similar operation for  $\Delta\tilde{a}_\mu(f)$  leads to

$$\Delta a_\mu(f) \simeq \frac{\alpha m_\mu^2}{4\pi^3 m_\phi^2} N(r_f) Q_f^2 \xi_\phi^\mu R_{i1}^\beta R_{i1}^\alpha \left[ \frac{(y + y^c)^2}{2} \mathcal{F}'_f(z) - \frac{yy^c}{z} \mathcal{F}_f(z) \right], \quad (25a)$$

$$\Delta\tilde{a}_\mu(f) \simeq \frac{\alpha m_\mu^2}{8\pi^3 m_\phi^2} N(r_f) Q_f^2 \xi_\phi^\mu R_{i1}^\beta R_{i2}^\beta (y^{c2} - y^2) \tilde{\mathcal{F}}'_f(z), \quad (25d)$$

where  $z = (m_l^2 + m_e^2)/2m_\phi^2$ . The above approximations are valid under the condition that  $m_l \approx m_e \gg \max(|y|, |y^c|) R_{i1}^\beta v$ .

Regarding the sign of  $\Delta\tilde{a}_\mu(f)$ , it is clear from (25d) that one can make it positive by choosing either  $|y|$  or  $|y^c|$  to be much larger than the other. Similarly, one can check that (25a) becomes positive for a pair of  $y$  and  $y^c$  with comparable magnitudes and appropriate signs. This is to be contrasted with vanilla NFC 2HDMs in which the sign of each Barr-Zee contribution is determined by the type of the Yukawa structure, the parity of the exchanged scalar, as well as whether the fermion in the loop is of up-type or down-type. For instance, the  $H$ - $\tau$  contribution to  $a_\mu$  in Type-II and -X is negative (see e.g. [35]). In a model extended with extra vector-like fermions, as shown above, there is enough freedom to engineer their gauge quantum numbers and Yukawa couplings so that their Barr-Zee contributions, mediated by  $H$  or  $A$ , have the desired sign.

We restrict ourselves hereafter to the  $i = 1$  case since coupling  $\Phi_2$  to the vector-like fermions would make an excessive modification to  $h \rightarrow \gamma\gamma$ . We plot the ratio  $\Delta\tilde{a}_\mu(f)/\tilde{c}_\gamma(f)$  in Fig. 4 without using mass insertion approximation. We set each of  $y$  and  $y^c$  to either 0 or 1 such that  $\Delta\tilde{a}_\mu(f) > 0$ . The overall size of  $y^{(c)}$  drops out of the ratio to a good approximation.

In Fig. 4(a) for scalar exchange, the curves resemble those in Fig. 3 although they lie lower than  $\Delta a_\mu(f)/c_\gamma(f)$  in the single-particle case, reproduced here as the thin solid curve from Fig. 3. This is because the contributions to each of  $\Delta a_\mu$  and  $c_\gamma$  from the two mass eigenstates of the vector-like fermions add up destructively. One can bring the  $i = 1$  curve

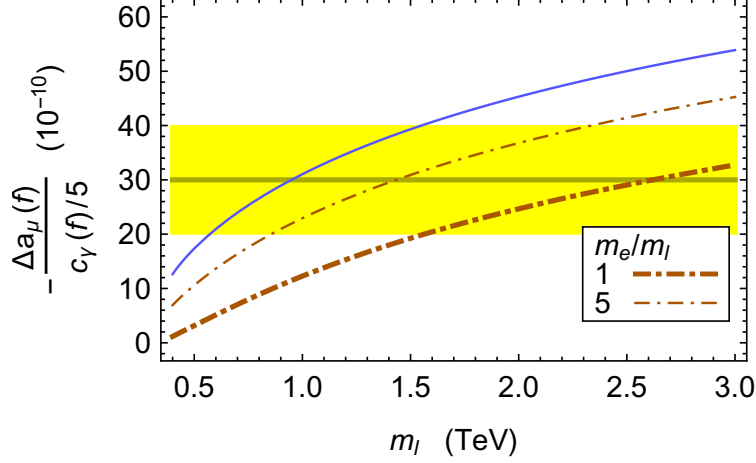
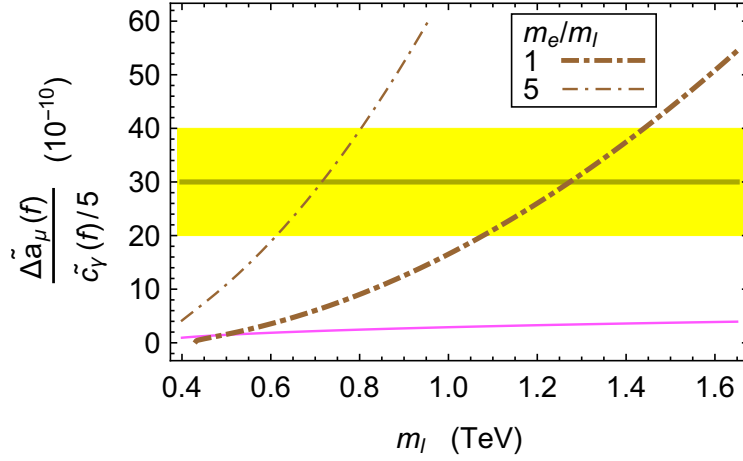
(a) Scalar exchange,  $\tan \beta = 10$ (b) Pseudoscalar exchange,  $\tan \beta = 1$ 

FIG. 4. (a)  $\Delta a_\mu(f)/c_\gamma(f)$  as a function of  $m_l$  with  $y = y^c$ . (b)  $\Delta \tilde{a}_\mu(f)/\tilde{c}_\gamma(f)$  as a function of  $m_l$  with  $y^c = 0$ . In both plots,  $m_e/m_l$  is fixed to 1 and 5 on the thick and the thin non-solid curves, respectively. The thin solid curves represent the values that would result from single-particle dominance.

closer to the single-particle result by decoupling the heavier state, as illustrated by the thin curve on which  $m_e/m_l = 5$ . The plot would remain the same under the interchange of  $m_l \leftrightarrow m_e$  in both the legend and the horizontal axis.

A similar plot for pseudoscalar exchange is shown in Fig. 4(b). Here, the curves have qualitatively different shapes from those in Fig. 4. Another difference is that the curves indicate much higher values of  $\Delta \tilde{a}_\mu(f)/\tilde{c}_\gamma(f)$  than the single-particle case. These differences can be traced to the ratio  $\tilde{\mathcal{F}}'_f(z)/[\tau^2 A'_a(\tau)]$  from (25d) and (24d), as opposed to  $\tilde{\mathcal{F}}_f(z)/A_a(\tau)$  appearing in (18d).

On the whole, Fig. 4 reveals a chance to explain  $\Delta a_\mu$  as well as the diphoton events by extending a 2HDM with vector-like fermions. Even though the curves therein look different from those in Fig. 3, it is still possible to find points with the best value of the ratio. Generalization to cases with multiple generations should be straightforward as long as those generations are arranged in such a way that their contributions add up constructively. To fit

the preferred size of  $\Delta a_\mu$ , the above particular model would in practice require a large number of generations, a large  $Q_f$ , and/or a large  $|y^{(c)}|$ , which might threaten the calculability of the model. In this respect, the scalar exchange contributions seem to be more promising. For  $m_l = m_e \sim 1$  TeV, one can account for  $\Delta a_\mu$  and  $c_\gamma$  by setting  $Q_f \sim 23$ ,  $y = y^c \sim 4$ , and  $\tan \beta \sim 25$ , which is still around the boundary of the calculability criteria from (3.25) of [6] with  $N_E = 3$ . Note that the pseudoscalar exchange contributions are suppressed in this case [see (24d) and (25d)]. The high value of  $\tan \beta$  was chosen to suppress the extra fermion loop contribution to  $h \rightarrow \gamma\gamma$  below  $2\sigma$  [26]. This would however be disfavoured by the limit on  $\Gamma_{\tau\tau}$  in a NFC model [see (20)]. As mentioned before, one might consider a non-NFC model in such a case.

To sum up, we attempted to predict generic contributions to the muon anomalous magnetic moment from a spin-zero particle  $\tilde{\phi}$  which we assume to produce the recently observed 750 GeV diphoton events. We evaluated the ratio of each component of  $\Delta a_\mu$  to the corresponding component of  $\tilde{c}_\gamma$ , the latter being the effective coupling which describes the decay  $\tilde{\phi} \rightarrow \gamma\gamma$ . Each ratio could be expressed as a simple function of the masses appearing in the loops thanks to the cancellation of the coupling constants common in the numerator and denominator.

We found: (1) if the 750 GeV resonance does not directly couple to the muon, the deviation  $\Delta a_\mu$  is difficult to explain; (2) if it has an enhanced coupling with the muon for instance as in the 2HDM Type-II or -X, then  $\Delta a_\mu$  can be easily accommodated. These observations bring us to the following enthralling interpretation of the latest LHC data: in conjunction with  $\Delta a_\mu$ , the 750 GeV diphoton excess suggests that the Standard Model Higgs sector should be extended by *more than*  $SU(2)_L$  singlet scalar fields.

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